

# Towards the Realistic Gluodynamics String. Perturbative Gluons' Contribution to the String Effective Action

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## Abstract

Perturbation theory in the nonperturbative QCD vacuum and the non-Abelian Stokes theorem, representing a Wilson loop in the  $SU(2)$  gluodynamics as an integral over all the orientations in colour space, are applied to derivation of the correction to the string effective action, obtained in Ref. 1. This correction is due to accounting in the lowest order of perturbation theory for the interaction of perturbative gluons with the string world sheet. It occurs that this interaction affects only the coupling constant of the rigidity term, while its contribution to the string tension of the Nambu-Goto term vanishes. The obtained correction to the rigid string coupling constant multiplicatively depends on the spin of the representation of the Wilson loop under consideration, the QCD coupling constant and a certain path integral, which includes the background Wilson average.

## 1. Introduction

Recently, a new approach to the gluodynamics string was suggested<sup>1</sup>. Within this approach one considers the Wilson average written through the non-Abelian Stokes theorem<sup>2,3</sup> and the cumulant expansion<sup>3,4</sup> as a statistical weight in the partition function of some effective string theory. The action of this string theory may be then obtained via expansion of the Wilson average in powers of two small parameters,  $(\alpha_s F_{\mu\nu}^a(0) F_{\mu\nu}^a(0))^{\frac{1}{2}} T_g^2$  and  $(\frac{T_g}{r})^2$ , where  $T_g \simeq 0,2 fm$  is the correlation length of the vacuum<sup>5,6</sup>, and  $r \simeq 1 fm$  is the size of the Wilson loop in the confining regime<sup>7</sup>. This yields the so-called curvature expansion for the gluodynamics string effective action. In Ref. 1 only the first nonvanishing terms of this expansion, corresponding to the lowest, second, order in the former parameter were accounted for, which corresponds to the so-called bilocal

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approximation<sup>3,8-11</sup>, and the expansion up to the terms of the third order in the latter parameter was elaborated out. The first two terms of this expansion read as follows

$$S_{biloc.} = \sigma \int d^2\xi \sqrt{g} + \frac{1}{\alpha_0} \int d^2\xi \sqrt{g} g^{ij} (\partial_i t_{\mu\nu}) (\partial_j t_{\mu\nu}), \quad (1)$$

where

$$\sigma = 4T_g^2 \int d^2z D(z^2) \quad (2)$$

is the string tension of the Nambu-Goto term, and

$$\frac{1}{\alpha_0} = \frac{1}{4} T_g^4 \int d^2z z^2 (2D_1(z^2) - D(z^2)) \quad (3)$$

is an inverse bare coupling constant of the rigidity term, while the terms of the third order in  $\left(\frac{T_g}{r}\right)^2$  contain higher derivatives of the induced metric  $g_{ij}$  and/or of the extrinsic curvature tensor  $t_{\mu\nu}$  w.r.t. world sheet coordinates, and we shall not quote them here (see Ref. 1 for the details)<sup>1</sup>. In Eqs. (2) and (3)  $D$  and  $D_1$  stand for two renormalization group invariant coefficient functions, parametrizing the gauge-invariant bilocal correlator of gluonic field strength tensors<sup>3,8-10</sup>, and it is worth noting that since the nonperturbative parts of these functions are related to each other as  $|D_1| \simeq \frac{1}{3}D$  according to lattice data<sup>6</sup>, the inverse rigid string coupling constant (3) is negative, which according to Ref. 12 agrees with the mechanism of confinement, based on the dual Meissner effect<sup>13</sup> (see discussion in Ref. 1). This result confirms that the Method of Vacuum Correlators, developed in Refs. 3 and 8-11 (see also Refs. therein), provides us with the consistent description of the confining gluodynamics vacuum. The approach suggested in Ref. 1 enables one to express all the coupling constants of the terms emerging in the string action in higher orders of the curvature expansion through the gauge-invariant correlators of gluonic field strength tensor only.

Notice also, that in Ref. 14 action (1) was applied to the derivation of the correction to the Hamiltonian of the QCD string with quarks, which was obtained in Ref. 15, arising due to the rigidity term, with the help of which a rigid string-induced term in the Hamiltonian of the relativistic quark model was then evaluated for the case of large masses of a quark and antiquark.

However, it should be emphasized that the curvature expansion describes only the pure non-perturbative content of the gluodynamics string theory. As it was explained in Ref. 10, in order to get the correct spectrum of the open bosonic string and the exponential growth of the multiplicity of string states, which is necessary for ensuring the property of duality of amplitudes, contained in the Veneziano formula, one must account for the perturbative gluons interacting with the string, which can be done in the framework of the perturbation theory in the nonperturbative QCD vacuum<sup>9</sup> (see discussion in Ref. 1).

In this letter, we shall take this interaction into account in the lowest order of perturbation theory and obtain the corresponding correction to the action (1). To this end, one needs to integrate over perturbative fluctuations in the expression for the Wilson average written through the non-Abelian Stokes theorem. This procedure is, however, looks rather difficult to be elaborated out in the case when one makes use of the version of the non-Abelian Stokes theorem, suggested in

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<sup>1</sup>From now on we shall use for the world sheet indices the letters from the middle of the Latin alphabet,  $i, j, k, \dots$ , in order not to confuse them with the colour indices  $a, b, c, \dots$ . We shall also keep the standard notation " $g$ " for the QCD coupling constant and would like to prevent the reader from confusing it with the determinant of the induced metric tensor, denoted by the same letter.

Refs. 2 and 3, due to the path-ordering, which is remained in the expression for the Wilson loop after rewriting it as a surface integral. In what follows, in order to get rid of it, we shall exploit another version of the non-Abelian Stokes theorem, which was proposed in Ref. 16, where the path-ordering was replaced by the integration over an auxiliary field from the  $SU(N_c)/[U(1)]^{N_c-1}$  coset space. For simplicity we shall consider the  $SU(2)$ -case, when this field is a unit three-vector  $\vec{n}$ , which characterizes the instant orientation in the colour space, and the non-Abelian Stokes theorem takes a remarkably simple form. Integration over perturbative fluctuations then yields the interaction of the elements of the string world sheet via the nonperturbative gluonic exchanges. In other words, we arrive at a theory of the open nonperturbative strings between dynamical world sheet elements, which provides us with the quantitative description of the intuitive picture, described in the previous paragraph. Finally, due to the short rangeness of this nonperturbative interaction between world sheet elements, it occurs possible to perform the  $\vec{n}$ -averaging and extract all the remnant dependence on the background fields in the form of the Wilson average standing under a certain path integral, so that the dependence on the world sheet elements decouples and may be evaluated explicitly, which yields a correction to the rigidity term, while the string tension of the Nambu-Goto term does not acquire any corrections and keeps its pure nonperturbative value (2). All the points, mentioned above, will be worked out in the next Section.

The main results of the letter are summarized in the Conclusion.

## 2. An Action of the Gluodynamics String Including Perturbative Gluons' Contributions

The statistical weight of the effective string theory, we are going to derive, is the Wilson average in the  $SU(2)$  gluodynamics  $\langle W(C) \rangle = \left\langle \text{tr } P \exp \left( ig \oint_C dx_\mu A_\mu^a t^a \right) \right\rangle$ , which after rewriting it as a surface integral by virtue of the non-Abelian Stokes theorem, suggested in Ref. 16, splitting the total field  $A_\mu^a$  into the background  $B_\mu^a$  and the perturbative fluctuations  $a_\mu^a$ ,  $A_\mu^a = B_\mu^a + a_\mu^a$ , and making use of the background field formalism<sup>9,17</sup>, takes the form

$$\begin{aligned} \langle W(C) \rangle = \int DB_\mu^a Da_\mu^a D\vec{n} \exp \left[ \int dx \left[ -\frac{1}{4} (F_{\mu\nu}^a)^2 + \frac{1}{2} a_\mu^a D_\lambda^{ab} D_\lambda^{bc} a_\mu^c + a_\nu^a D_\mu^{ab} F_{\mu\nu}^b + \right. \right. \\ \left. \left. + \frac{iJ}{2} \int d\sigma_{\mu\nu} n^a \left( -g (F_{\mu\nu}^a + 2D_\mu^{ab} a_\nu^b) + \varepsilon^{abc} (D_\mu \vec{n})^b (D_\nu \vec{n})^c \right) \right] \right], \end{aligned} \quad (4)$$

where  $F_{\mu\nu}^a = \partial_\mu B_\nu^a - \partial_\nu B_\mu^a + g\varepsilon^{abc} B_\mu^b B_\nu^c$  is a strength tensor of the background field,  $D_\mu^{ab} = \delta^{ab} \partial_\mu - g\varepsilon^{abc} B_\mu^c$  is the corresponding covariant derivative, and  $J = \frac{1}{2}, 1, \frac{3}{2}, \dots$  is the spin of the representation of the Wilson loop under consideration. In what follows we shall be interested only in the effects of the lowest order of perturbation theory, so that on the R.H.S. of Eq. (4) we have replaced the full covariant derivative which should stand in the last term (the so-called Wess-Zumino term) by the background one and omitted the ghost term, the term, describing the interaction of two perturbative gluons with the background field strength tensor and with the string world sheet, and the terms which describe the interaction of three and four perturbative gluons. Integration over the perturbative fluctuations in Eq. (4) is Gaussian and yields

$$\langle W(C) \rangle = \int DB_\mu^a D\vec{n} \exp \left[ -\frac{1}{4} \int dx (F_{\mu\nu}^a)^2 + \frac{iJ}{2} \int d\sigma_{\mu\nu} \varepsilon^{abc} n^a (D_\mu \vec{n})^b (D_\nu \vec{n})^c \right].$$

$$\begin{aligned}
& \cdot \exp\left(-\frac{iJg}{2} \int d\sigma_{\mu\nu} n^a F_{\mu\nu}^a\right) \exp\left[-\frac{1}{2} \int dx dy \left[D_{\mu}^{ba} \left(iJg T_{\mu\nu}^a(x) + F_{\mu\nu}^a(x)\right)\right] \int_0^{\infty} ds \int (Dz)_{xy} e^{-\int_0^s \frac{\dot{z}^2}{4} d\lambda}\right. \\
& \left. \cdot \left[P \exp\left(ig \int_0^s d\lambda \dot{z}_{\alpha} B_{\alpha}\right)\right]^{bc} \left[D_{\rho}^{cd} \left(iJg T_{\rho\nu}^d(y) + F_{\rho\nu}^d(y)\right)\right]\right], \tag{5}
\end{aligned}$$

where  $T_{\mu\nu}^a(x) \equiv \int d^2\xi \varepsilon^{ij} (\partial_i x_{\mu}(\xi)) (\partial_j x_{\nu}(\xi)) n^a(x(\xi)) \delta(x - x(\xi))$  is the colour vorticity tensor current. Making use of the formula<sup>4</sup>

$$\langle e^P Q \rangle = \left( \exp\left(\sum_{n=1}^{\infty} \frac{1}{n!} \langle\langle P^n \rangle\rangle\right) \right) \left( \langle Q \rangle + \sum_{k=1}^{\infty} \frac{1}{k!} \langle\langle P^k Q \rangle\rangle \right),$$

where  $P$  and  $Q$  stand for two statistically dependent commuting quantities, and  $\langle\langle \dots \rangle\rangle$  denotes the so-called cumulants, i.e. irreducible correlators, we get in the lowest order of the cumulant expansion the following correction to the string effective action (1)

$$\begin{aligned}
\Delta S = & -\frac{J^2 g^2}{2} \int dx dy \left\langle \left( D_{\mu}^{ba} T_{\mu\nu}^a(x) \right) \int_0^{\infty} ds \int (Dz)_{xy} e^{-\int_0^s \frac{\dot{z}^2}{4} d\lambda} \right. \\
& \left. \cdot \left[ P \exp\left(ig \int_0^s d\lambda \dot{z}_{\alpha} B_{\alpha}\right) \right]^{bc} D_{\rho}^{cd} T_{\rho\nu}^d(y) \right\rangle_{\vec{n}, B_{\mu}^a}, \tag{6}
\end{aligned}$$

where

$$\begin{aligned}
\langle \dots \rangle_{\vec{n}} & \equiv \int D\vec{n} (\dots) \exp\left(\frac{iJ}{2} \int d\sigma_{\mu\nu} \varepsilon^{abc} n^a (D_{\mu} \vec{n})^b (D_{\nu} \vec{n})^c\right), \\
\langle \dots \rangle_{B_{\mu}^a} & \equiv \int DB_{\mu}^a (\dots) \exp\left(-\frac{1}{4} \int dx (F_{\mu\nu}^a)^2\right),
\end{aligned}$$

and during the derivation of Eq. (6) we have omitted the term

$$\exp\left[-\frac{1}{2} \int dx dy \left(D_{\mu}^{ba} F_{\mu\nu}^a(x)\right) \int_0^{\infty} ds \int (Dz)_{xy} e^{-\int_0^s \frac{\dot{z}^2}{4} d\lambda} \left[ P \exp\left(ig \int_0^s d\lambda \dot{z}_{\alpha} B_{\alpha}\right) \right]^{bc} D_{\rho}^{cd} F_{\rho\nu}^d(y) \right]$$

on the R.H.S. of Eq. (5), which does not yield any contributions to the string action due to the lack of coupling with the world sheet and therefore may be absorbed into the measure  $DB_{\mu}^a$ . Integrating in Eq. (6) by parts, we arrive in the lowest order of perturbation theory at the following formula

$$\begin{aligned}
\Delta S = & -\frac{J^2 g^2}{2} \int d^2\xi \int d^2\xi' \varepsilon^{ij} \varepsilon^{kl} (\partial_i x_{\mu}) (\partial_j x_{\nu}) (\partial_k x'_{\rho}) (\partial_l x'_{\nu}) \cdot \\
& \cdot \left\langle \left\langle n^b(x) n^c(x') \right\rangle_{\vec{n}} \frac{\partial^2}{\partial x_{\mu} \partial x'_{\rho}} \int_0^{\infty} ds \int (Dz)_{xx'} e^{-\int_0^s \frac{\dot{z}^2}{4} d\lambda} \left[ P \exp\left(ig \int_0^s d\lambda \dot{z}_{\alpha} B_{\alpha}\right) \right]^{bc} \right\rangle_{B_{\mu}^a}, \tag{7}
\end{aligned}$$

where  $x_\mu \equiv x_\mu(\xi)$ , and  $x'_\mu \equiv x_\mu(\xi')$ .

Since  $e^{-\int_0^s \frac{\dot{z}^2}{4} d\lambda} P \exp \left( ig \int_0^s d\lambda \dot{z}_\alpha B_\alpha \right)$  is the statistical weight of a perturbative gluon, propagating from the point  $x'$  to the point  $x$  along the trajectory  $z_\alpha$  during the proper time  $s$ , it is the region where  $s$  is small, which mainly contributes to the path integral on the R.H.S. of Eq. (7). This means that the dominant contribution to  $\Delta S$  comes from those  $x_\mu$ 's and  $x'_\mu$ 's, which are very close to each other, which is in the line with the curvature expansion, where  $|x' - x| \leq T_g \ll r$ . Within this approximation, one gets

$$\langle n^b(x) n^c(x') \rangle_{\vec{n}} \simeq \frac{\delta^{bc}}{3} \int D\vec{n} \exp \left( \frac{iJ}{2} \int d\sigma_{\mu\nu} \varepsilon^{def} n^d (D_\mu \vec{n})^e (D_\nu \vec{n})^f \right). \quad (8)$$

It is worth noting, that the integral on the R.H.S. of Eq. (8) is a functional of the world sheet as a whole (it is independent of  $x_\mu(\xi)$ ), and therefore may be also referred to the measure  $DB_\mu^a$ .

Hence, as it was announced in the Introduction, we see that expression (7) for the correction to the string effective action (1) due to the perturbative gluons takes the form of the interaction of two elements of the world sheet,  $d\sigma_{\mu\nu}(\xi)$  and  $d\sigma_{\rho\nu}(\xi')$ , via the nonperturbative gluonic string.

In order to extract explicitly the dependence on the points  $x$  and  $x'$  from the functional integral standing on the R.H.S. of Eq. (7), let us pass to the integration over the trajectories  $u_\mu(\lambda) = z_\mu(\lambda) + \frac{\lambda}{s}(x' - x)_\mu - x'_\mu$ , which yields

$$\Delta S = J^2 g^2 \int d^2 \xi \int d^2 \xi' \varepsilon^{ij} \varepsilon^{kl} (\partial_i x_\mu) (\partial_j x_\nu) (\partial_k x'_\rho) (\partial_l x'_\nu) \frac{\partial^2}{\partial x_\mu \partial x'_\rho} \int_0^\infty ds e^{-\frac{(x-x')^2}{4s}} \int (Du)_{00} e^{-\int_0^s \frac{\dot{u}^2}{4} d\lambda} \cdot \left\langle \text{tr} P \exp \left[ ig \int_0^s d\lambda \left( \frac{x-x'}{s} + \dot{u} \right)_\alpha B_\alpha \left( u + x' + \frac{\lambda}{s}(x-x') \right) \right] \right\rangle_{B_\mu^a}, \quad (9)$$

and from now on we shall absorb the inessential constant factors into the measure  $DB_\mu^a$ .

The Wilson loop standing on the R.H.S. of Eq. (9) may be expanded as follows

$$\begin{aligned} \text{tr} P \exp \left[ ig \int_0^s d\lambda \left( \frac{x-x'}{s} + \dot{u} \right)_\alpha B_\alpha \left( u + x' + \frac{\lambda}{s}(x-x') \right) \right] &= \text{tr} P \exp \left( ig \int_0^s d\lambda \dot{u}_\alpha B_\alpha(u) \right) + \\ &+ ig \text{tr} \int_0^s d\lambda \left[ P \exp \left( ig \int_0^\lambda d\lambda' \dot{u}_\alpha B_\alpha(u(\lambda')) \right) \right] \left[ \frac{(x-x')_\beta}{s} B_\beta(u(\lambda)) + \right. \\ &\left. + \left( x' + \frac{\lambda}{s}(x-x') \right)_\beta \dot{u}_\gamma(\lambda) (\partial_\beta B_\gamma(u(\lambda))) \right] \left[ P \exp \left( ig \int_\lambda^s d\lambda'' \dot{u}_\zeta B_\zeta(u(\lambda'')) \right) \right] + O(g^2), \end{aligned}$$

and since we are working in the lowest order of perturbation theory, all the terms of this expansion except for the first one will be omitted below, so that Eq. (9) yields

$$\Delta S = J^2 g^2 \int d^2 \xi \int d^2 \xi' \varepsilon^{ij} \varepsilon^{kl} (\partial_i x_\mu) (\partial_j x_\nu) (\partial_k x'_\rho) (\partial_l x'_\nu).$$

$$\int_0^\infty \frac{ds}{s} \left( \frac{(x-x')_\mu (x-x')_\rho}{2s} - \delta_{\mu\rho} \right) e^{-\frac{(x-x')^2}{4s}} \Phi(s), \quad (10)$$

where

$$\Phi(s) \equiv \int (Du)_{00} e^{-\int_0^s \frac{\dot{u}^2}{4} d\lambda} \left\langle \text{tr } P \exp \left( ig \int_0^s d\lambda \dot{u}_\alpha B_\alpha(u) \right) \right\rangle_{B_\mu^a}. \quad (11)$$

Finally, in order to get the desirable correction to the action (1), we shall expand the R.H.S. of Eq. (10) in powers of  $\frac{s}{r^2}$  (according to the discussion in the paragraph before Eq. (8)), keeping in this expansion terms not higher in the derivatives w.r.t. world sheet coordinates than the rigidity, which corresponds to the expansion up to the second order in the parameter  $\frac{s}{r^2}$ . Omitting the full derivative terms of the form  $\int d^2\xi \sqrt{g} R$ , where  $R$  is a scalar curvature of the world sheet, we get analogously to Ref. 1 the following values of the integrals standing on the R.H.S. of Eq. (10)

$$\begin{aligned} & \int d^2\xi \int d^2\xi' \varepsilon^{ij} \varepsilon^{kl} (\partial_i x_\mu) (\partial_j x_\nu) (\partial_k x'_\rho) (\partial_l x'_\nu) (x-x')_\mu (x-x')_\rho \int_0^\infty \frac{ds}{s^2} e^{-\frac{(x-x')^2}{4s}} \Phi(s) = \\ & = 4\pi \int_0^\infty ds \Phi(s) \left( 4 \int d^2\xi \sqrt{g} - 3s \int d^2\xi \sqrt{g} g^{ij} (\partial_i t_{\mu\nu}) (\partial_j t_{\mu\nu}) \right) \end{aligned} \quad (12)$$

and

$$\begin{aligned} & \int d^2\xi \int d^2\xi' \varepsilon^{ij} \varepsilon^{kl} (\partial_i x_\mu) (\partial_j x_\nu) (\partial_k x'_\mu) (\partial_l x'_\nu) \int_0^\infty \frac{ds}{s} e^{-\frac{(x-x')^2}{4s}} \Phi(s) = \\ & = 2\pi \int_0^\infty ds \Phi(s) \left( 4 \int d^2\xi \sqrt{g} - s \int d^2\xi \sqrt{g} g^{ij} (\partial_i t_{\mu\nu}) (\partial_j t_{\mu\nu}) \right). \end{aligned} \quad (13)$$

Combining together Eqs. (12) and (13), we arrive at the following correction to the effective action (1) due to the accounting for the perturbative gluons in the lowest order of perturbation theory

$$\Delta S = \left( \Delta \frac{1}{\alpha_0} \right) \int d^2\xi \sqrt{g} g^{ij} (\partial_i t_{\mu\nu}) (\partial_j t_{\mu\nu}), \quad (14)$$

where

$$\left( \Delta \frac{1}{\alpha_0} \right) = J^2 g^2 \int_0^\infty ds s \Phi(s). \quad (15)$$

Notice, that as it was already pointed out in the Introduction, perturbative gluons do not change the value of the string tension (2) of the Nambu-Goto term and affect only the coupling constant of the rigidity term. Since this correction (15) to the nonperturbative rigid string coupling constant (3) is a pure perturbative effect, its sign is unimportant for the explanation of confinement in terms of the dual Meissner effect (see discussion in the Introduction), which allowed us to refer the constant factor in Eq. (15) to the measure  $DB_\mu^a$  in Eq. (11). The nontrivial content of this

correction emerges due to the path integral defined by Eq. (11), which includes the background Wilson average.

### 3. Conclusion

In this letter we have applied perturbation theory in the nonperturbative QCD vacuum<sup>9</sup> and the non-Abelian Stokes theorem, which represents a Wilson loop in the  $SU(2)$  gluodynamics as an integral over all the orientations in colour space<sup>16</sup> to the derivation of the correction to string effective action (1), found in Ref. 1, which emerges due to accounting for the interaction of perturbative gluons with the string world sheet in the lowest order of perturbation theory. This correction is given by formulae (11), (14) and (15) and affects only the rigidity term, while the string tension of the Nambu-Goto term keeps its pure nonperturbative value (2). Perturbative correction (15) to the inverse coupling constant of the rigidity term contains the dependence on the background fields in the form of the background Wilson average standing under a certain path integral (11).

We have also demonstrated that perturbative fluctuations, when being taken into account, lead to the interaction defined by the R.H.S. of Eq. (7) between elements of the string world sheet by virtue of nonperturbative gluonic strings, which agrees with the qualitative scenario of excitation of the gluodynamics string by the perturbative gluons, suggested in Refs. 1 and 10.

However, it is still remains unclear whether perturbative gluons may yield cancellation of the conformal anomaly in  $D = 4$  rather than in  $D = 26$ , as it takes place for the ordinary bosonic string theory<sup>18,19</sup>, and the solution of the problem of crumpling for the rigidity term<sup>19,20</sup>. These problems will be treated in the next publications.

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